The dilaton-dominated supersymmetry breaking scenario in the context of the non-minimal supersymmetric model

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ABSTRACT

The phenomenological consequences of the dilaton-type soft supersymmetry breaking terms in the context of the next to minimal supersymmetric standard model are investigated. We always find a very low top quark mass. As a consequence such string vacua are excluded by recent experimental results. The viability of the solution of the μ term through the introduction of a gauge singlet field is also briefly discussed.

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The Standard Model remains unscathed in its confrontation with experiment. Furthermore, the compatibility of the model with the recent CDF observation [1], is very encouraging and enhances our belief that we are on the right path towards the realization of the unification program.

However, in order to talk about the unification of all the forces we have to see the standard model as the effective limit of a more fundamental theory. Today the only theoretical framework which is a strong candidate for the unification of all the interactions, including gravity, is heterotic string theory. Supersymmetry which is necessary in taming the radiative corrections to the Higgs boson mass, the only still unseen particle of the standard model, is naturally embedded in string theory. However, supersymmetry has to be softly-broken at an energy of order of the electroweak scale in order to solve the hierarchy problem mentioned above. In the Minimal Supersymmetric extension of the Standard Model (MSSM) the terms responsible for the breaking are customarily parametrized in terms of four universal parameters, $M_{1/2}$, A, m_0, B, the universal gaugino mass, the trilinear scalar terms associated with the trilinear couplings in the superpotential, the scalar masses, and the B term associated with the higgs-doublet mixing term in the superpotential. In string theory these parameters are in principle, predicted but a definite answer is at present lacking due to the fact that the supersymmetry breaking mechanism in the theory is not well understood.

Nevertheless progress has taken place at the theoretical level, and therefore we believe that it is an appropriate time to use all the present existing knowledge to study string inspired scenarios and confront them with current experimental data. In this way phenomenological criteria may help us select the correct string vacuum state among the *plethora* of equivalent string vacua that appear at the level of string pertubation theory, the majority of them leading to unacceptable phenomenology.

The process of supersymmetry breaking has to have a non-pertubative origin since it is well known that SUSY is preserved order by order in pertubation theory. However, very little is known about non-pertubative effects in string theory, particularly in the four-dimensional case. This has led the authors of [3], to parametrize the effect of SUSY-breaking by the VEVs of the F-terms of the dilaton (S) and the moduli (T_m) [2] chiral superfields generically present in large classes of four-dimensional supersymmetric heterotic strings . In a way, if supersymmetry breaking is triggered by these fields (i.e., $\langle F_S \rangle \neq 0$ or $\langle F_T \rangle \neq 0$), this would be a rather generic prediction of string theory.

There are various possible scenarios for supersymmetry breaking which are obtained in this model independent way. To discriminate among these we consider a simplified expression for the scalar masses

$$m_i^2 = m_{3/2}^2 (1 + n_i \cos^2 \theta) \tag{1}$$

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[2]In string pertubation theory, both the moduli and the dilaton are exact flat directions of the effective potential, leaving their VEVs undetermined. In ref. [3] this pertubative degeneracy is assumed to be completely lifted by the non-pertubative dynamics and VEVs for moduli and dilaton to be induced.

with $\tan\theta=\langle F_S\rangle/\langle F_T\rangle$ [3]. Here $m_{3/2}$ is the gravitino mass and the n_i are the modular weights of the respective matter field. There are two ways in which one can obtain universal scalar masses, as strongly desired phenomenologically to avoid large flavor-changing-neutral currents (FCNCs) [4]: (i) setting $\theta = \pi/2$, that is $\langle F_S\rangle\gg\langle F_T\rangle$; or (ii) in a model where all n_i are the same, as occurs for $Z_2\times Z_2$ orbifolds [3] and free fermionic constructions [5]. In the first scenario supersymmetry breaking is triggered by the dilaton F- term and if the vanishing of the cosmological constant is imposed as a constraint, this results in a universal scenario for the soft parameters involved [3]. In particular we have:

$$A = -M_{1/2}, \quad m_0 = \frac{1}{\sqrt{3}}M_{1/2} \tag{2}$$

As one can see the four-dimensional soft-parameter space is then effectively two-dimensional. In the strict-dilaton case the B term is no longer an independent parameter but is given by

$$B = \frac{2}{\sqrt{3}}M_{1/2} = 2m_0 \tag{3}$$

The above universal scenario leads to a natural supression of FCNC. The dilaton dominated scenario has been studied in the context of the minimal supersymmetric model [8] and in the context of the flipped SU(5) [9]. However, in both cases the necessary Higgs mixing term is provided by a bilinear mass term. The relevant term in the superpotential is of the form

$$W_\mu = \mu H_1.H_2. \tag{4}$$

The associated soft-breaking term in the scalar potential will have the form

$$BH_1.H_2 \tag{5}$$

It is well known that, in order to get appropriate $SU(2)_L\times$U(1) breaking, the μ parameter has to be of the same order of magnitude as the SUSY-breaking soft terms. This is in general unexpected since the μ-term is a supersymmetric term whereas the other soft terms are originated after supersymmetry breaking. This is called "the μ problem", the reason why μ should be of the order of the soft terms.

A very attractive solution to the problem [10] is to add an extra singlet superfield N which couples with the two Higgs doublets superfields. The resulting superpotential will contain besides the usual standard model terms the following term.

$$W_N = \lambda N H_1.H_2 + \frac{1}{3}kN^3 \tag{6}$$

If, the gauge singlet acquires a vacuum expectation value $\langle N\rangle$ then the role of μ is played by $\lambda\langle N\rangle$ and the role of B is played by A_λ. The resulting model is the so called nonminimal supersymmetric standard model. We note that such a superpotential which contains only trilinear couplings emerges naturally in superstrings.

It is also very interesting to point out that in general the soft terms computed in a general class of string models are complex. The resulting phases, are quite constrained by limits on the electric dipole moment of the neutron (EDMN), since they give large one-loop contributions to this CP-violating quantity. In the case of the MSSM, there are three classes of phases associated with the parameters A, $M_{1/2}$, B and μ which are candidates for CP violation. These are given by

$$
\begin{aligned}
\phi_A &= arg(A_{ijk}/\lambda_{ijk}), \\
\phi_B &= arg(B/\mu), \\
\phi_C &= arg(M_a)
\end{aligned}
\tag{7}
$$

where M_a are the masses of the gaugino fields associated with the three gauge group factors of the MSSM and λ_{ijk} are the trilinear Yukawa couplings of the chiral superfields of the theory and A_{ijk} the usual trilinear soft terms. However, it has been shown [6]that after a redefinition there are only two CP violating phases

$$
\phi_A = arg(AM^*_{1/2}), \quad \phi_B = arg(BM^*_{1/2})
\tag{8}
$$

These phases are constrained as we mentioned by the electric dipole moment of the neutron. One has in fact

$$
\phi_A, \phi_B \leq 10^{-3}
\tag{9}
$$

for sparticle masses around few hundreds GeV. The EDMN receives important contributions from both ϕ_A and ϕ_B phases. In the dilaton-dominated scenario in the MSSM it has been shown that the phases for M_a and A coincide and $\phi_A \to 0$ [3]. However, ϕ_B is in general large and can be sufficiently supressed only if further assumptions about the origin of SUSY-breaking are made [3, 7].

In the case of the nonminimal model, the dilaton-dominated scenario leads to a natural supression of the EDMN. Indeed, if the μ problem is solved by the addition of a singlet N, the role of B is played by a trilinear coupling A_λ whose phase is aligned with that of the gauginos and naturally $\phi_A = \phi_{A_\lambda} = 0$. Thus the nonminimal model with supersymmetry breaking by the VEV of the dilaton field is a well motivated scenario that deserves further study. However, as we shall see, phenomenological considerations exclude the model.

A very attractive feature of spontaneously broken effective supergravities is the radiatively induced breaking of the electroweak gauge symmetry [13].For the study of electroweak symmetry breaking we use the one-loop effective potential

$$
V = V_0 + \Delta V_1
\tag{10}
$$

where the radiative corrections to the scalar potential

$$
\Delta V_1 = \frac{1}{64\pi^2} \text{Str} \left[\mathcal{M}^4 \left(\ln \frac{\mathcal{M}^2}{Q^2} - \frac{3}{2} \right) \right]
\tag{11}
$$

depend on the Higgs fields through the tree-level squared-mass matrix $\mathcal{M}^2$. The supertrace in equation (11) is given by

$$\mathrm{Str} f(\mathcal{M}^2) = \sum (-1)^{2J_i} (2J_i + 1) f(m_i^2) \tag{12}$$

where m_i^2 denotes the field-dependent mass eigenvalue of the ith particle of spin J_i. In the above expression for V, all parameters of the theory are running parameters, functions of the renormalization point Q. The use of the one-loop effective potential guarantees the scale-independence of the solutions [12]. The tree level potential V_0 contains the standard F- and D- terms, and the following soft-supersymmety breaking terms:

$$(M_1\lambda_1\lambda_1 + M_2\lambda_2\lambda_2 + M_3\lambda_3\lambda_3 + h_t A_t Q.H_2 U_R^c + \lambda A_\lambda H_1.H_2 N + \frac{1}{3}kA_k N^3) + h.c.$$
$$+ m_1^2|H_1|^2 + m_2^2|H_2|^2 + m_N^2|N|^2 + m_Q^2|Q|^2 + m_{U^c}^2|U^c|^2 + \dots \tag{13}$$

where λ_1, λ_2, and λ_3 denote the gauginos of the $U(1)_Y$, SU(2) and SU(3) gauge groups, respectively.

For the calculation of radiative corrections in (11) we take into account the corrections due to top, and stop loops. This approximation is correct as long as $\tan\beta < m_t/m_b$, where m_t and m_b are the top and bottom quark masses respectively and $\tan\beta$ is the ratio of the VEVs of the two Higgs doublets

$$\tan\beta = \langle H_2\rangle / \langle H_1\rangle \tag{14}$$

It has been shown that supersymmetry prevents the spontaneous breaking of CP [11], so that the vevs of the fields H_1, H_2 and N are of the form

$$\langle H_1\rangle = \begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \langle H_2\rangle = \begin{pmatrix} 0 \\ v_2 \end{pmatrix}, \langle N\rangle = x \tag{15}$$

with v_1, v_2, x real.

The equations for extrema of the full scalar potential in the directions (15) in field space read

$$v_1[m_1^2 + \lambda^2(v_2^2 + x^2) + \frac{1}{4}(g_1^2 + g_2^2)(v_1^2 - v_2^2)] + \lambda v_2 x(kx + A_\lambda) + \frac{1}{2}\partial\Delta V_1/\partial v_1 = 0, \tag{16}$$

$$v_2[m_2^2 + \lambda^2(v_1^2 + x^2) + \frac{1}{4}(g_1^2 + g_2^2)(v_2^2 - v_1^2)] + \lambda v_1 x(kx + A_\lambda) + \frac{1}{2}\partial\Delta V_1/\partial v_2 = 0, \tag{17}$$

$$x[m_N^2 + \lambda^2(v_1^2 + v_2^2) + 2k^2x^2 + 2\lambda k v_1 v_2 + kA_k x] + \lambda A_\lambda v_1 v_2 + \frac{1}{2}\partial\Delta V_1/\partial x = 0 \tag{18}$$

We now describe our numerical procedure. At the string unification scale the following relations hold

$$g_1 = g_2 = g_3 \equiv g_U$$

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$$M_1 = M_2 = M_3 \equiv M_{1/2}$$
$$h_t = h_U^t, \lambda = \lambda_U, k = k_U \tag{19}$$
$$m_i^2 \equiv m_0^2 = \frac{1}{3} M_{1/2}^2, \quad i = 1, 2, N, Q, U^c$$
$$A_t = A_\lambda = A_k \equiv A = -M_{1/2}$$

We scan the parameter space of the model as follows:First we choose a set of initial values for the parameters,h_U^t, λ_U, k_U, $M_{1/2}$ at the scale of O(10^{16})GeV. Then by using the well known set of RGE at the one-loop order we evolve the relevant parameters down to the electroweak scale of O(100)GeV and insert their numerical values into the minimization equations $(16) - (18)$. The latter are highly non-linear algebric equations with independent variables the three vacuum expectation values. In order to solve them we use numerical routines from the NAG library which iteratively converge to a solution. We start with an initial guess for the three vacuum expectation values and the routines employed quickly converge to the solution for v_1, v_2, x. We seek solutions where three nonzero vacuum expectation values develop. Our solutions are nontrivially constrained, by the physical requirement that the correct mass for the Z^0 boson must be reproduced. Furthermore, in order to guarantee a global minimum, we demand that all the physical Higgs bosons have positive mass squared and that the vacuum expectation value of the Higgs potential is negative and therefore energetically preferable over the symmetric minimum. An essential check of the correctness of the minimization of the one-loop effective potential is the appearance of the charged and neutral massless Goldstone bosons, which indicate that the charged and neutral Higgs field dependence has been included properly in the $\mathcal{M}$ matrices appearing in (10). The running top quark mass obtained in our procedure is related to the experimentally observable pole mass by [14]

$$m_t^{pole} = m_t(m_t) \left[1 + \frac{4}{3} \frac{\alpha_s(m_t)}{\pi} + K_t \left(\frac{\alpha_s(m_t)}{\pi} \right)^2 \right] \tag{20}$$

where

$$K_t = 16.11 - 1.04 \sum_{i=1}^{5} \left(1 - \frac{M_i}{M_t} \right) \tag{21}$$

where M_i, i=1,...,5, represent the masses of the five lighter quarks. We solve (20) by using the nonlinear equation solution routines described above for the minimization of the effective potential.

The experimental constraints that we impose are summarized in table 1. We also require that $m_{\tilde{\nu}}^2 > 0$ to avoid a $\Delta L \neq 0$ vacuum [10] and that $m_t^{pole} > 131 GeV$ [2].

Particle	Experimental Limit (GeV)
gluino	120
squark,slepton	45
chargino	45
neutralino	20
light higgs	60

Table 1: Experimental constraints

As it turned out the model failed to satisfy all the constraints imposed from radiative electroweak breaking and the experiment. The results of our research are summarized in table 2. In particular we list values of the relevant parameters which reproduce the correct Z^0 boson mass, $m_{\tilde{\nu}}^2 > 0$ and gluino masses above $120 GeV$. The single entry for m_t^{pole} in table 2 indicates the maximum value of the physical top quark mass obtained in our study. For our notation in table 2 see Ellis et al. in ref. [10].

As one can see from the solutions of the renormalization group the values of the trilinear Yukawa couplings of the model that give correct electroweak breaking are very constrained. In particular the top Yukawa coupling is always small implying a very low top quark mass in conflict with published experimental results [2]. Higher values of the top Yukawa coupling require very small gluino masses in order that the constraint of correct electroweak breaking is satisfied. Also higher values for h_U^t tend to generate VEVs for the sneutrino field. We do not list in table 2 lower values than 0.16 for the h_U^t coupling.The lightest Higgs mass eigenstate, m_{S_1}, is always below the threshold of 60 GeV. Also, from the study of the vacuum expectation value of the gauge singlet Higgs field we conclude that the generated μ term is only a few GeV, implying therefore a light chargino mass,$m_{\chi_1^+}$, sometimes below the current experimental limit. At this point we must mention that it is possible to generate a different set of VEVs, which give $m_t \sim 130 GeV$, by starting with higher values for the top and λ Yukawas. The feature of the latter set is the large vacuum expectation value of the gauge singlet of O(1)TeV. However, this set of VEVs lead to an unstable vacuum since $\langle V_{Higgs}(1 - loop)\rangle > 0$ and the Higgs mass squared eigenvalues are not all positive. This is in agreement with the work in ref.[15] where an upper bound for the λ_U Yukawa has been obtained, $\lambda_U < 0.55$. In our case the values of λ_U which generates the large VEV for the gauge singlet are such that $\lambda_U \geq 0.65$ for $h_U^t \geq 0.35$. All the above facts make the nonminimal model with dilaton-dominated soft-terms excluded by experiment. It is worthwhile to mention that the no-scale scenario with the supersymmetry breaking driven by the F-terms of the moduli fields [3]is also not a viable scenario in the context of supersymmetric models with a gauge singlet [15].

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[3] The dilaton-dominated scenario is a no-scale scenario driven by the F-terms of the dilaton superfield.

Parameter	$|k_U| = 0.01$	$|k_U| = 0.1$	$|k_U| = 1.0$
h_U^t	0.16-0.182	0.16-0.184	0.16-0.189
λ_U	0.16-0.35	0.19-0.35	0.34-0.57
h_t	0.47-0.52	0.47-0.53	0.46-0.54
λ	0.2-0.403	0.24-0.405	0.34-0.5
$|k|$	0.01	0.082-0.093	0.43-0.49
$m_t(\mathrm{GeV})$	71-91	71 -90	70-91
m_t^{pole}	97	96	96
$\tan\beta$	1.42-7.6	1.41-4.51	1.4-3.1
r	0.09-0.52	0.21-0.52	0.34-0.5
$m_{\tilde{g}}(\mathrm{GeV})$	120-192	121-200	120-218
$m_{\tilde{\nu}}(\mathrm{GeV})$	2.3-37	3.12-39	9.2-51
$m_{\tilde{e_R}}(\mathrm{GeV})$	40-68	41-66	41-69
$m_{\chi_1^+}(\mathrm{GeV})$	12-67	28-68	43-71
$m_{\chi_1^0}(\mathrm{GeV})$	23-35	23-38	9-39
$m_{H^\pm}(\mathrm{GeV})$	77-89	78-89	82-98
$m_{P_1}(\mathrm{GeV})$	6.6-11	29-36	62-103
$m_{P_2}(\mathrm{GeV})$	62-108	74-104	104-114
$m_{S_1}(\mathrm{GeV})$	10-39	21-39	12-44
$m_{S_2}(\mathrm{GeV})$	49-92	59-88	89-97
$m_{S_3}(\mathrm{GeV})$	95-108	95-108	108-134

Table 2: Typical parameter values that emerge from the renormalization group analysis which are consistent with correct electroweak breaking

Thus the supersymmetric models with a gauge singlet Higgs field seem difficult to reconcile with universal no-scale scenarios. We remind the reader that no-scale supergravity is the infrared limit of superstring theory [16]. Alternative mechanisms for the generation of the μ term exist in the context of string theory [3]. From the phenomenological point of view string inspired univeral no-scale supergravity models with two Higgs doublets are preferable to models with the extra gauge singlet Higgs field. However, the latter with more general soft-supersymmetry breaking terms can lead to acceptable phenomenology though is difficult to have definite predictions for the particle spectrum due to the high-dimensionality of the parameter space [15]. On the other hand, if the problem of the μ term in string theory is solved by the introduction of a gauge singlet superfield this might lead to a departure from the universality of the soft SUSY breaking terms. [4].

Acknowledgments

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[4] Non universal scenarios for the soft-terms emerge naturally on orbifold constructions [3].

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